

ELECTRON-CAPTURE AND LOW-MASS IRON-CORE-COLLAPSE SUPERNOVAE:
NEW NEUTRINO-RADIATION-HYDRODYNAMICS SIMULATIONSDAVID RADICE^{1,2}, ADAM BURROWS², DAVID VARTANYAN², M. AARON SKINNER³, AND JOSHUA C. DOLENCE⁴*Draft version February 15, 2017*

ABSTRACT

We present results from new 1D (spherical) and 2D (axisymmetric) long-term FORNAX simulations of electron-capture (EC) and low-mass iron-core-collapse supernovae. We consider six progenitor models: the $8.8 M_{\odot}$ (zero-age main sequence [ZAMS]) prototype ECSN progenitor from [Nomoto \(1984, 1987\)](#); two ECSN-like iron core progenitors, with ZAMS masses of $8.1 M_{\odot}$, and $9.6 M_{\odot}$ ([Müller et al. 2012a, 2013](#)), with $10^{-4} Z_{\odot}$ and zero metallicity, respectively; and the 9-, 10-, and $11 M_{\odot}$ (ZAMS) progenitors from [Sukhbold et al. \(2016\)](#). We confirm that the ECSN and ECSN-like progenitors explode easily even in 1D with explosion energies of up to a quarter of a Bethe ($1 B \equiv 10^{51}$ erg). We also find that ECSNe are a viable mechanism for the production of very low-mass neutron stars. However, the 9-, 10-, and $11 M_{\odot}$ progenitors do not explode in 1D and are not even necessarily easier to explode than higher-mass progenitor stars in 2D. We discuss the consequences of our findings for currently proposed “explodability” criteria. Furthermore, we study the effect of perturbations and of changes to the microphysics. We find that relatively small changes can result in qualitatively different outcomes, even in 1D, for models sufficiently close to the explosion threshold. Finally, we revisit the impact of convection below the protoneutron star (PNS) surface. Leveraging the 1D explosions, we analyzed, for the first time, 1D and 2D evolutions of PNSs with the same boundary conditions. We find that the impact of PNS convection has been underestimated in previous studies and results in an increase by a factor ~ 2 of the neutrino luminosity for all species starting from roughly half a second after bounce.

Keywords: Stars: supernovae: general

1. INTRODUCTION

The formation of a massive iron core at the end of the evolution of stars with zero-age main-sequence (ZAMS) masses larger than $\sim 12 M_{\odot}$ is a robust prediction of stellar evolution theory. These stars undergo core-collapse once their cores reach the Chandrasekhar mass and may explode as core-collapse supernovae (CCSNe).

The fate of less massive stars in the ZAMS range $\sim 8 M_{\odot}$ to $\sim 12 M_{\odot}$ is less clear. Depending on the initial mass, the ultimate fate could be to form massive white dwarfs, to form iron cores as regular massive stars, or to explode as electron-capture supernovae (ECSNe) before forming an iron core (e.g. [Nomoto 1984, 1987](#); [Jones et al. 2013](#); [Doherty et al. 2015](#); [Woosley & Heger 2015](#)). It has also been suggested that these stars might undergo violent flashes and power unusual transients before their deaths ([Woosley & Heger 2015](#); [Jones et al. 2016](#)).

ECSNe and low-mass iron-core CCSNe with similar features are expected to occur in a relatively narrow range of ZAMS masses. However, they might account for a significant fraction of gravitational-collapse SNe given that the initial mass function of stars drops rapidly towards high masses. These progenitors have compact cores with tenuous envelopes, which result in a steep drop of the accretion rate after core bounce. This, in turn, triggers early explosions, even under the assump-

tion of spherical symmetry ([Kitaura et al. 2006](#); [Janka et al. 2008](#); [Burrows et al. 2007](#); [Fischer et al. 2010](#); [Janka et al. 2012](#)). ECSNe and ECSN-like CCSNe are expected to be underenergetic and possibly underluminous and to have small ^{56}Ni yields and peculiar nucleosynthetic abundances (e.g. [Nomoto et al. 1982](#); [Kitaura et al. 2006](#); [Janka et al. 2008](#); [Hoffman et al. 2008](#); [Wanajo et al. 2011](#); [Melson et al. 2015b](#); [Wanajo et al. 2017](#)).

[Kitaura et al. \(2006\)](#) suggested that ECSN-like events might explain a subclass of Type-II SNe with unusually low luminosities ([Pastorello et al. 2004](#); [Spiro et al. 2014](#)). An ECSN has also been invoked to explain SN 1054 and the associated Crab remnant ([Nomoto et al. 1982](#); [Takahashi et al. 2013](#); [Smith 2013](#); [Tominaga et al. 2013](#)). According to historical records, SN 1054 was not underluminous. However, SN-1054 was likely underenergetic, with an explosion energy around 10^{50} erg, as indicated by the low-mass of the Crab nebula’s filaments and their relatively low expansion velocity, as well as by the small inferred ^{56}Ni yield ([Müller 2016](#), and references therein). On the basis of measured isotopic abundance anomalies, it has also been suggested that a low-mass CCSN might have been the trigger that started the formation of our Solar System ([Banerjee et al. 2016](#)).

ECSNe and ECSN-like progenitors have attracted significant interest in the CCSN mechanism community due to their “explodability” and the fact that they allow for self-consistent studies also in 1D (with the assumption of spherical symmetry). [Hillebrandt et al. \(1984\)](#) performed the first 1D simulations of the collapse, bounce, and explosion of the original $n8.8$ ECSN progenitor of [Nomoto \(1984, 1987\)](#), using an approximate gray neutrino transport scheme. They found an energetic explo-

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sion by the prompt shock mechanism with an energy of $\sim 2 \cdot 10^{51}$ erg. Subsequent studies, with modern neutrino interactions and multi-group transport in 1D and 2D, performed by Kitaura et al. (2006); Janka et al. (2008); Burrows et al. (2007); Fischer et al. (2010), found much weaker explosions ($\sim 10^{50}$ erg) powered by the delayed neutrino mechanism. Müller et al. (2012a) considered an $8.1 M_\odot$ (ZAMS) progenitor with $Z = 10^{-4}$, u8.1, which formed an iron core but had a stellar structure very similar to the n8.8 progenitor, and found a similarly early explosion. Another iron-core progenitor, the zero-metallicity $9.6 M_\odot$ progenitor, z9.6, was considered and determined to have a qualitatively similar outcome by Janka et al. (2012); Müller et al. (2013); Müller & Janka (2014) in 2D and by Melson et al. (2015b) in 3D. Very recently, Wanajo et al. (2017) presented new nucleosynthetic calculations for the n8.8, u8.1, and z9.6 models, as well as a summary of their associated explosion characteristics with the COCONUT-VERTEX code.

The low-mass, but otherwise “canonical,” $11.2 M_\odot$ progenitor from Woosley et al. (2002) and the $12.0 M_\odot$ progenitor from Woosley & Heger (2007) have been considered by several groups, (e.g. Buras et al. 2006; Takiwaki et al. 2012; Müller et al. 2012b; Bruenn et al. 2013, 2016; Dolence et al. 2015; Müller 2015; Summa et al. 2016; O’Connor & Couch 2015; Burrows et al. 2016; Nagakura et al. 2017). While the $11.2 M_\odot$ progenitor has a post-bounce evolution that is qualitatively similar to the u8.1 progenitor (Müller et al. 2012a), the $12.0 M_\odot$ progenitor either explodes very late (Summa et al. 2016; Burrows et al. 2016) or not at all (O’Connor & Couch 2015; Burrows et al. 2016), with the exception of the simulations by Bruenn et al. (2013, 2016). This challenges the commonly held idea that low-mass iron-core progenitors should explode easily and similarly to an ECSN.

With the goal of characterizing the different explosions of ECSNe, ECSN-like CCSN, and canonical, but low-mass CCSNe, we present 1D and 2D FORNAX simulations of the collapse, bounce, and subsequent evolution of six progenitor models. We consider the ECSN n8.8 progenitor from Nomoto (1984, 1987), the ECSN-like u8.1 and z9.6 progenitors from Müller et al. (2012a) and Müller et al. (2013), and the 9-, 10-, and $11 M_\odot$ solar-metallicity iron-core collapse progenitors from Sukhbold et al. (2016). We show that low-mass CCSNe always fail to explode in 1D and evolve in a qualitatively different way compared to ECSNe and ECSN-like CCSNe. We also discuss the impact of the dimensionality, of pre-supernova perturbations, and of changes in the microphysics. In particular, for the latter, we focus on the effects of many-body corrections to the axial-vector term in the neutrino-nucleon scattering rate recently proposed by Horowitz et al. (2016).

The rest of this paper is organized as follows. First, in Section 2, we give an overview of the simulation setup and of the properties of the progenitor models. We discuss the qualitative outcome of our simulations in Section 3, while a more quantitative account of the energetics of the explosions is given in Section 4. We discuss the properties of the neutrino radiation in Section 5. Section 6 is dedicated to the properties and evolution of the remnant protoneutron stars (PNSs). Finally, we summarize and discuss our results in Section 7.

2. PROGENITORS AND SETUP

As previously discussed, we consider six progenitor models, which we label as n8.8 (Nomoto 1984, 1987), u8.1 (Müller et al. 2012a), z9.6 (Müller et al. 2013), and 9.0-, 10.0-, $11.0 M_\odot$ (Sukhbold et al. 2016). All of the progenitors have been evolved up to the point of core-collapse, defined as the time their radial infall velocity has reached $\sim 1000 \text{ km s}^{-1}$. Their structure (density and exterior binding energies) are shown in Fig. 1. All of these progenitors have relatively compact cores and loosely bound envelopes. The values of the compactness parameter $\xi_{2.5}$ (O’Connor & Ott 2011, 2013) computed from the progenitor models are given in Table 2. They range from $\simeq 7.6 \cdot 10^{-5}$, for the z9.6 progenitor, to $\simeq 7.7 \cdot 10^{-3}$, for the $11.0 M_\odot$ progenitor. $\xi_{2.5}$ cannot be computed for the n8.8 progenitor since it only extends to $\simeq 1.32 M_\odot$. Note that, for numerical reasons, we modify the n8.8 progenitor for $\rho \leq 10^4 \text{ g cm}^{-3}$ with the addition of a constant temperature envelope with $\rho \propto r^{-2}$. We verified, by changing the density of the envelope over 3 orders of magnitude, that the explosion energy of the n8.8 progenitor is not sensitive to it, although the shock propagation speed is obviously affected.

All of the progenitors are non-rotating and have been evolved in 1D. There is currently significant interest in the possible impact of pre-supernova asphericities on the development of explosions, which several authors have shown to be beneficial and, in some cases, determinant of the outcome of the evolution (Couch & Ott 2013; Müller & Janka 2015; Abdikamalov et al. 2016; Takahashi et al. 2016). The first progenitor models evolved in 3D shortly before core collapse have also recently become available (Couch et al. 2015; Müller et al. 2016). Here, we also consider the impact of perturbations on the 9.0-, 10.0-, and $11.0 M_\odot$ progenitors using the approach introduced by Müller & Janka (2015), which we briefly describe. We introduce velocity perturbations in three regions $r_i^- \leq r \leq r_i^+$, $i = 1, 2, 3$, obeying the divergence-free condition $\nabla \cdot (\rho \delta \mathbf{v}_i) = 0$. These are generated as

$$\delta \mathbf{v}_i = \begin{cases} \frac{C_i}{\rho} \nabla \times \Psi_i, & r_i^- \leq r \leq r_i^+, \\ 0, & \text{otherwise;} \end{cases} \quad (1)$$

where

$$\Psi_i = \mathbf{e}_\phi \frac{\sqrt{\sin \theta}}{r} \sin \left(n_i \pi \frac{r - r_i^-}{r_i^+ - r_i^-} \right) Y_{\ell_i, 1}(\theta, 0) \quad (2)$$

and n_i , ℓ_i are the number of convective cells in the radial and angular directions, respectively. Finally, C_i is tuned to achieve a given maximum perturbation amplitude. The parameters we use are given in Table 1.

We evolve these progenitors from the onset of collapse with the neutrino-radiation-hydrodynamics FORNAX code (Skinner et al. 2015; Burrows et al. 2016; Dolence et al. 2017 in prep) and a setup similar to the one in Burrows et al. (2016).

FORNAX solves for the transport of neutrinos using a multi-dimensional moment scheme with an analytic closure for the 2nd and 3rd moments (Shibata et al. 2011; Murchikova et al. 2017). Similar moment methods have also been recently adopted for the CCSN problem by other groups (e.g. O’Connor 2015; Just et al. 2015; O’Connor & Couch 2015; Roberts et al. 2016). The moment equations are solved using a 2nd-order finite-volume

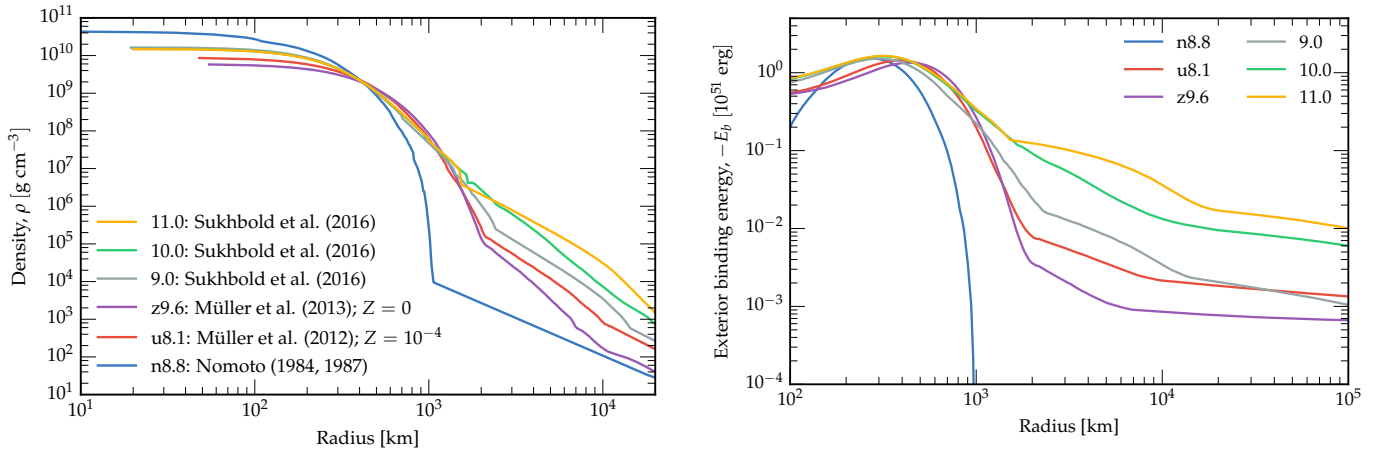


Figure 1. Progenitor models: density profiles in g cm^{-3} (left panel) and binding energies in Bethes ($1 \text{ B} = 10^{51} \text{ erg}$; right panel). The envelope binding energy is computed as the total energy exterior to a given radius. Note that, for numerical reasons, we modified the n8.8 progenitor with the addition of a thicker envelope (see the main text for details). Progenitors that successfully explode in 1D (n8.8, u8.1, and z9.6) have steeper density profiles and smaller binding energies than low-mass progenitors that do not explode in 1D.

Table 1
Details of the setup for the perturbed models.

Prog.	r_1^- [km]	r_1^+ [km]	ℓ_1	n_1	δv_1 [10^8 cm s^{-1}]	r_2^- [km]	r_2^+ [km]	ℓ_2	n_2	δv_2 [10^8 cm s^{-1}]	r_3^- [km]	r_3^+ [km]	ℓ_3	n_3	δv_3 [10^8 cm s^{-1}]
9.0	850	1,350	3	1	1.0	2,500	9,750	1	1	0.3	15,000	450,000	1	1	0.3
10.0	480	560	10	1	1.0	1,150	1,800	1	3	1.0	2,500	3,400	5	1	0.5
11.0	450	520	11	1	0.5	1,100	1,400	7	1	1.2	1,550	11,000	1	1	0.5

scheme with the HLLE approximate Riemann solver (Einfeldt 1988), modified as in Audit et al. (2002) and O’Connor (2015) to reduce the numerical dissipation in the diffusive limit. FORNAX separately evolves electron neutrinos ν_e and anti-electron neutrinos $\bar{\nu}_e$, while heavy-lepton neutrinos ν_μ , ν_τ , and the respective anti-particles are lumped together as a single species, which we denote as “ ν_μ .” The energy spectra of neutrinos are resolved using 20 logarithmically-spaced energy groups extending to 300 MeV for electron neutrinos and to 100 MeV for anti-electron and heavy-lepton neutrinos. Neutrino-matter interactions are treated as discussed in Burrows et al. (2016). We perform three variants of each simulation. “Baseline” includes all the neutrino-matter interaction discussed in Burrows et al. (2016), with the exception of the many-body corrections to the neutrino-nucleon scattering cross-sections of Horowitz et al. (2016). “NoINS” includes the same reactions as Baseline with the exception of inelastic scattering on nucleons. “Horowitz” includes all of the Baseline reactions as well as the many-body corrections of Horowitz et al. (2016). For the runs with perturbations, we use the Baseline physics setup.

The hydrodynamic equations are solved using a high-resolution shock-capturing scheme with 3rd-order reconstruction and the HLLC approximate Riemann solver (Toro et al. 1994). The details of the numerical schemes are discussed in (Skinner et al. 2015; Burrows et al. 2016; Dolence et al. 2017 in prep). For the simulations presented here, we use a spherical grid with 678 points extending up to 20,000 km. The grid has a constant spacing Δr of 0.5 km for $r \lesssim 10$ km and then smoothly transitions to a logarithmically spaced grid with $\Delta r/r \simeq 0.01$ for $r \gtrsim 100$ km. For the 2D simulations, we use 256

angular zones with angular resolution smoothly varying between $\simeq 0.95^\circ$ at the poles and $\simeq 0.64^\circ$ at the equator. The angular grid is also progressively derefined towards the center to avoid an excessively restrictive CFL condition in the angular direction.

We adopt the Lattimer-Swesty equation of state with nuclear compressibility parameter 220 MeV (Lattimer & Swesty 1991) and treat gravity in the monopole approximation using a general-relativistic potential, following Marek et al. (2006).

Finally, we carry out all 2D simulations until the maximum shock radius exceeds 19,000 km, or until the explosion is deemed unsuccessful. We continue 1D simulations until the time when the minimum electron fraction in the PNS becomes equal to the minimum value of the equation of state table (0.035).

3. OVERALL DYNAMICS

A first glance of our results can be gained from Fig. 2, which shows the average shock radii for all progenitors in 1D and 2D with both the Baseline and Horowitz setups. As in previous works by others (Kitaura et al. 2006; Janka et al. 2008; Burrows et al. 2007; Fischer et al. 2010; Müller et al. 2012a; Janka et al. 2012; Müller et al. 2013; Melson et al. 2015b), we find early explosions for the n8.8, z9.6, and u8.1 progenitors. The same progenitors also explode in 1D spherical symmetry, although the u8.1 fails to explode in 1D if inelastic neutrino-nucleon scatterings are neglected, suggesting that its explosion is only marginal.

None of the progenitors from Sukhbold et al. (2016) explode in self-consistent 1D simulations. Somewhat surprisingly, the $10.0\text{-}M_\odot$ progenitor fails to explode also in

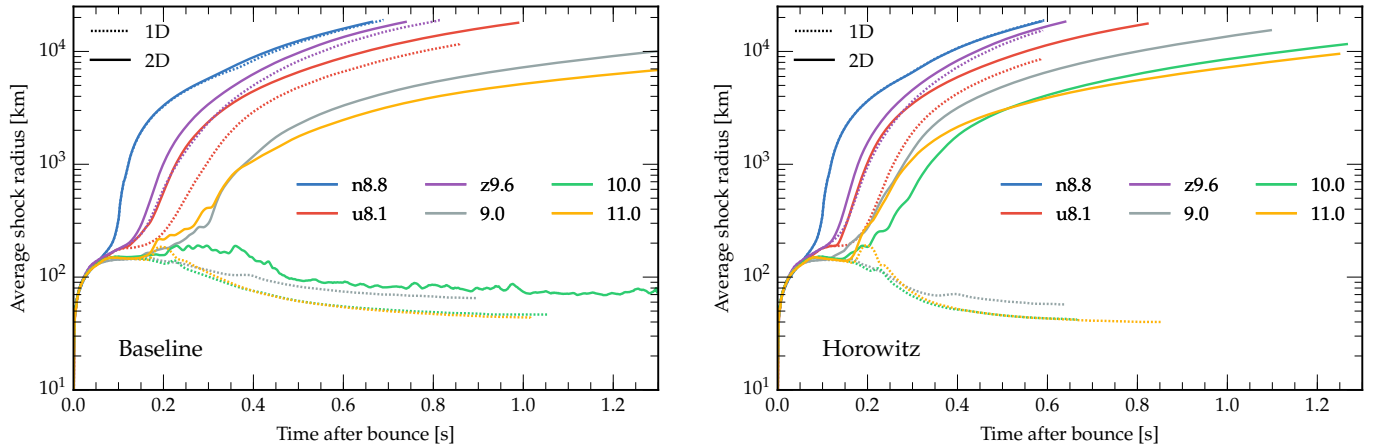


Figure 2. Average shock radius (km) tracks for all progenitors in 1D and 2D with our Baseline setup (*left panel*) and with the inclusion of many-body corrections (*right panel*). The curves are smoothed using a running average with a 5-ms window. With many-body corrections, the $10.0\text{-}M_{\odot}$ progenitor also explodes in 2D and the 9.0- and $11.0\text{-}M_{\odot}$ 2D explosions become more robust. None of the progenitors from Sukhbold et al. (2016) explode in 1D, even with many-body corrections.

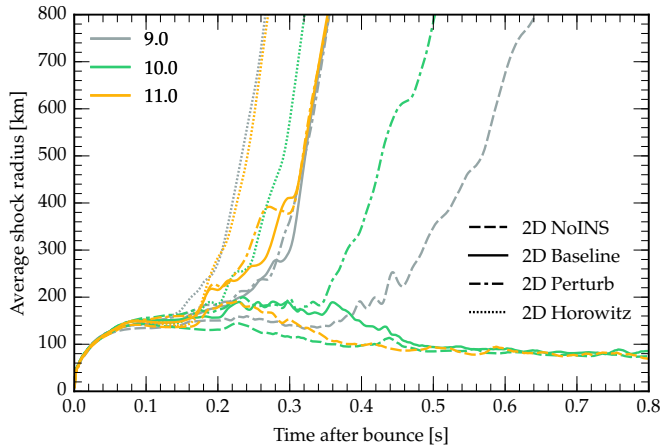


Figure 3. Impact of perturbations and changes to the microphysics on the average shock radius (km) of the 9.0- , 10.0- , and $11.0\text{-}M_{\odot}$ progenitors. The curves are smoothed using a running average with a 5-ms window. Both the inclusion of perturbations and changes to the microphysics (NoINS, Baseline, and Horowitz setups) can trigger explosions in models that otherwise fail.

2D with the Baseline setup within the simulation time ($\simeq 1.3$ s after bounce). We remark that in Burrows et al. (2016) we found successful explosions for the 20- and $25\text{-}M_{\odot}$ progenitors from Woosley & Heger (2007) using the same microphysical treatment. However, we did not find explosions for the 12- and $15\text{-}M_{\odot}$ progenitors without the inclusion of many-body corrections to the neutrino-nucleon scattering cross-sections. All progenitors explode successfully in 2D with the Horowitz setup.

To visualize the different outcomes of the progenitors from Sukhbold et al. (2016) in 2D, we highlight their early post-bounce average shock radii in Fig. 3. The $9.0\text{-}M_{\odot}$ and $11.0\text{-}M_{\odot}$ models have delayed explosions at ~ 0.3 s after bounce (Baseline setup) or ~ 0.2 s after bounce (Horowitz). The $10.0\text{-}M_{\odot}$ model also fails to explode. This is not surprising, given that the $10.0\text{-}M_{\odot}$ Baseline also fails, and in the light of the discussions in Müller et al. (2012b) and Burrows et al. (2016). These authors showed that the inclusion of inelastic scattering of heavy-lepton neutrinos near the ν_e and $\bar{\nu}_e$ neutrinospheres boosts the average energy of the former and re-

sults in an improved energy deposition rate. This effect makes the NoINS setup less explosive. The $11.0\text{-}M_{\odot}$ also fails to explode. The $9.0\text{-}M_{\odot}$ NoINS explodes much later than either the $9.0\text{-}M_{\odot}$ Horowitz or the $9.0\text{-}M_{\odot}$ Baseline, when the Si/O-O interface is accreted. The inclusion of perturbations does not affect the outcome of the 9.0- and $11.0\text{-}M_{\odot}$ progenitors, but is sufficient to trigger the explosion of the $10.0\text{-}M_{\odot}$.

The inclusion of perturbations results only in modest changes for the 9.0- and $11.0\text{-}M_{\odot}$ progenitors, which already explode at early times with the Baseline setup. Note that the amplitude of the initial perturbations are on the upper end of what could be considered as realistic, with turbulent velocities reaching $\sim 1000\text{ km s}^{-1}$ (Table 1). Despite the large perturbation amplitudes, the shock expansion is triggered only slightly earlier for the 9.0- and $11.0\text{-}M_{\odot}$ models. Then, starting from ~ 0.3 s after bounce, the shock trajectories for the 9.0- and $11.0\text{-}M_{\odot}$ Baseline and 9.0- and $11.0\text{-}M_{\odot}$ Perturb are very similar (Fig. 3). The $10.0\text{-}M_{\odot}$ Perturb is the only model for which we find perturbations to yield a qualitative change to the evolution and trigger a weak explosion ~ 0.4 s after bounce.

Some insight into the reason for the different evolutions can be gained from the analysis of the accretion rate history of the progenitors, as recently suggested by Suwa et al. (2016) and Müller (2016). Fig. 4 shows the accretion rate at 500 km for all progenitors in 1D, with the Baseline setup. Since the 1D progenitors from Sukhbold et al. (2016) fail to explode, these are “intrinsic” accretion rates, not affected by the explosion. The accretion rates for the n8.8, z9.6, and u8.1 are unaffected by the explosion up to the point where they are shown (afterwards, they become negative as the inflow turns into an outflow). It is easily seen that progenitors exploding in 1D have a steep decline of the accretion rate at very early times, which sets them apart from “normal” massive stars, as also pointed out by Müller (2016). We find that the $10.0\text{-}M_{\odot}$ progenitor accretion rate is significantly higher than the 9.0- and $11.0\text{-}M_{\odot}$ models during the critical phase when the other two start exploding. The sudden growth of the accretion rate of the $10.0\text{-}M_{\odot}$ progenitor around ~ 0.2 s after bounce is due to a den-

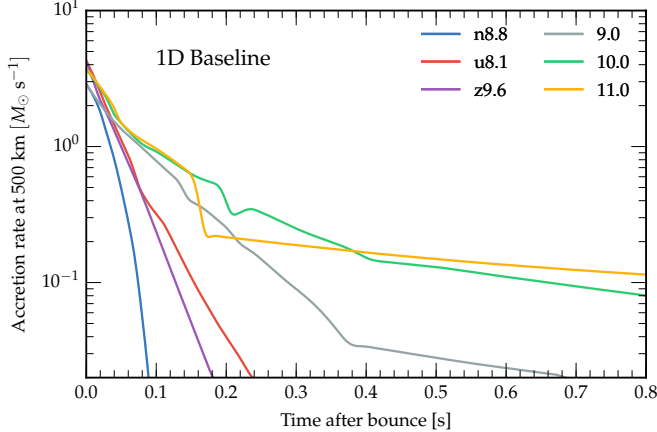


Figure 4. Accretion rates in $M_{\odot} \text{ s}^{-1}$ at 500 km for our Baseline 1D setup. The curves are smoothed using a running average with a 5-ms window. Successful 1D explosions require steep drops in the accretion rate at early times, when the neutrino luminosities are still large.

sity inversion present in the original progenitor profile (Fig. 1).

Figures 5 and 6 show snapshots of the entropy and of the pressure contrast, defined following Fernández & Thompson (2009) to be $r |\nabla p|/p$, for selected progenitors at 0.15 s, 0.2 s, and 0.4 s after bounce. In all our models, neutrino-driven convection is seeded in the region behind the shock by perturbations induced by the dendritic grid and develops rapidly following bounce (*i.e.*, starting from ~ 0.15 s after bounce). Large-scale shock-sloshing motions, possibly due to the standing accretion shock instability (SASI) (Blondin et al. 2003; Foglizzo et al. 2007), are only present at late times in models that fail to explode (the 10.0-Baseline and 10.0-NoINS, and the 11.0-NoINS), after the shock has receded to less than 100 km in radius. Convection is the dominant instability at early times and for all exploding models, due to the low post-bounce accretion rates (Foglizzo et al. 2006; Burrows et al. 2012; Murphy et al. 2013; Müller et al. 2012a; Ott et al. 2013; Couch & O’Connor 2014; Fernández et al. 2014; Abdikamalov et al. 2015).

That said, convection appears to have a different role in the onset of the explosion of the n8.8 and z9.6 compared to the 9.0-, 10.0-, and 11.0- M_{\odot} progenitors (for the setups which result in explosions). For the n8.8 and the z9.6, the shock starts expanding rapidly already after about one convective overturn, before the high-entropy plumes are able to reach it (Figs. 5 and 6). The onset of these two explosions is essentially spherically symmetric and does not seem to be significantly aided by convection, although there are small differences between 1D and 2D for the z9.6 visible in Fig. 2. We point out that Melson et al. (2015b) found a significantly more delayed explosion in 1D for the z9.6 as compared to their 2D and 3D simulations: compare our 1D and 2D shock radius evolution in Fig. 2 to theirs in Fig. 3 of Melson et al. (2015b).

Different from the n8.8 and z9.6, the 9.0-, 10.0-, 11.0- M_{\odot} progenitors, explode only after a few convective overturns with the emergence of one or more large plumes that succeed in pushing the shock to a sufficiently large radius to trigger a run-away expansion. This behavior is commonly observed in 2D CCSN simulations (*e.g.*,

Fernández et al. 2014). The u8.1 is intermediate in that the convective plumes manage to reach and deform the shock prior to the onset of the explosion (Fig. 6). However, also for the u8.1, the shock expansion appears to be triggered in a spherically-symmetric way by the drop of the accretion rate. Subsequently, the shock expands rapidly and becomes almost spherical.

After the explosion sets in, the shock and the material immediately behind it expand almost self-similarly, with roughly constant velocities. Behind them, we observe the emergence of a higher-entropy neutrino-driven wind with dynamics similar to the one reported by Burrows et al. (1995). For the n8.8, z9.6, and u8.1 progenitors the wind is quasi-spherical, it produces weak shocks visible in the pressure-contrast visualizations in Figs. 5 and 6, and drives Rayleigh-Taylor instabilities as it pushes on the slower, heavier material above. In the case of asymmetric explosions, the wind is typically confined in $\sim 90^\circ$ wedges along the axis where it drives the inflation of a large bubble, while fall-back accretion continues along the equator. We observed qualitatively similar features for the 20- M_{\odot} progenitor from Woosley & Heger (2007) in (Burrows et al. 2016; Vartanyan et al. 2017 in prep). We caution the reader, however, that the degree of asymmetry in the 9.0-, 10.0-, and 11.0- M_{\odot} progenitor explosions is likely to be artificially magnified by the assumption of axisymmetry, and we speculate that the neutrino-driven wind will be closer to spherical in full-3D simulations.

As a representative example of an ECSN-like explosion, we show in Figs. 7 and 8 a summary of the evolution of the z9.6 progenitor in 1D and 2D, evolved with the Baseline setup. Despite their relatively similar average shock trajectories and the fact that the z9.6 shock remains nearly spherical also in 2D, there are substantial differences between the 1D and 2D explosions. In particular, the entropy in the 1D model exceeds that of the 2D model, while the velocity of the neutrino-driven wind with the 2D model exceeds that in the 1D model, roughly by a factor of two. This is partly due to the tendency of 1D models to create low-density regions that become over-heated by neutrinos. More importantly, starting from ~ 0.3 s after bounce, the 2D simulation shows significantly larger neutrino luminosities (see Sec. 6 for a detailed account). This results in stronger neutrino-driven winds, with increased velocities (see Fig. 8) and correspondingly smaller expansion timescales.

The evolution of the n8.8 and u8.1 models is qualitatively and quantitatively similar to the z9.6. In contrast, the 9.0-, 10.0-, and 11.0- M_{\odot} progenitors evolve in a qualitatively different way. Fig. 9 shows the evolution in 2D of the 9.0 progenitor with the Baseline setup, which we take as representative of regular (successful) CCSNe from a low-mass progenitor. The explosion of the 9.0 follows ~ 100 ms of shock stagnation and is triggered at the time when the Si/O interface is accreted. As already mentioned, the explosion is asymmetric. It is also marginal, with small velocities immediately below the shock and a degree of sustained fallback at the end of the simulation.

4. EXPLOSION ENERGETICS

We estimate explosion energies using a fixed-volume energy analysis similar to that in Bruenn et al. (2016). We consider the region $r \geq 100$ km, and we compute

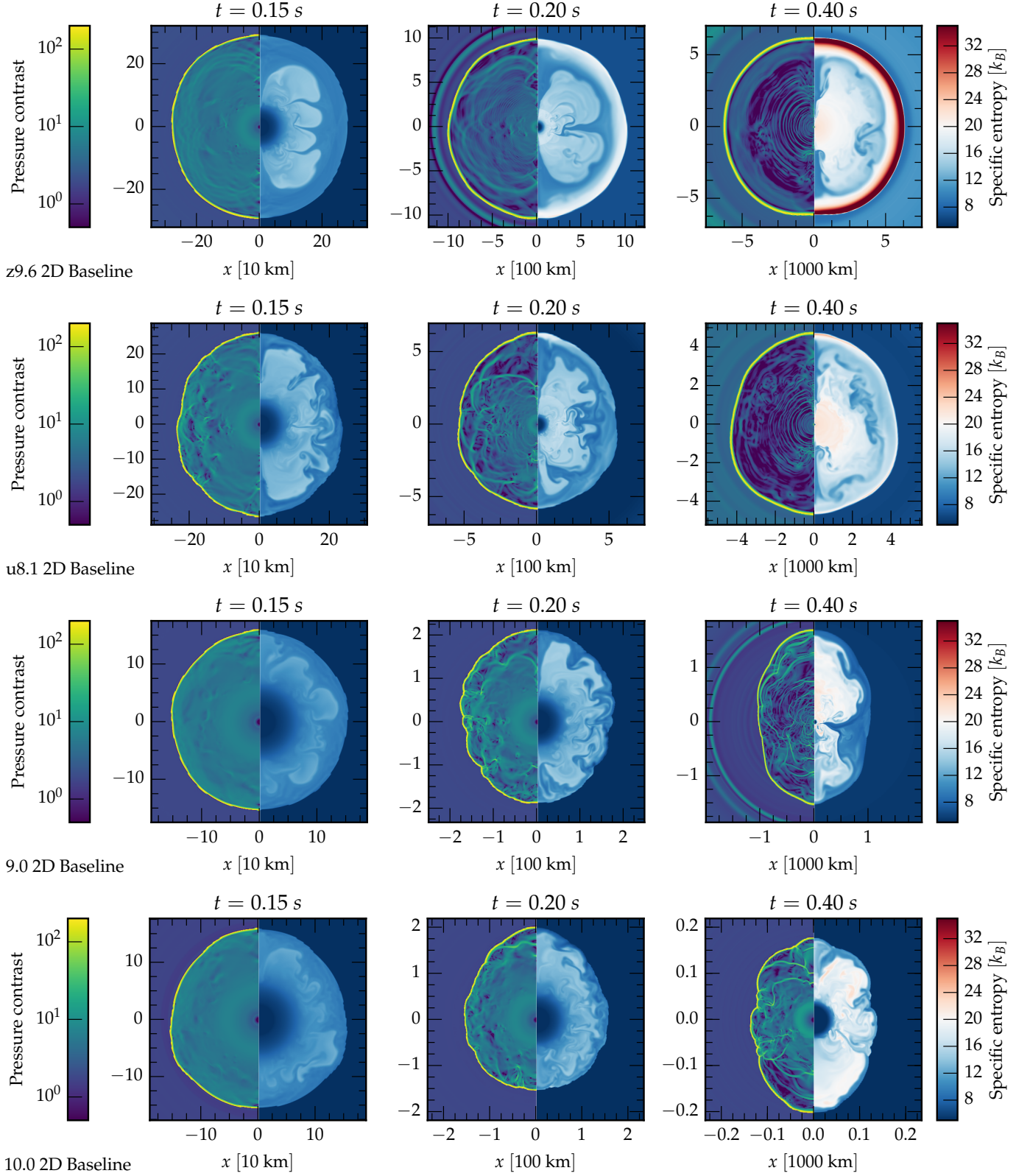


Figure 5. Entropy per baryon in k_B and pressure contrast ($r \cdot |\nabla p|/p$) profiles for the z9.6, u8.1, 9.0, and 10.0 progenitors evolved with the Baseline setup at three representative times. Note the different spatial scales. The ring-like structure visible in the pressure contrast in some panels are compositional shells. All models develop convection around ~ 0.15 s after bounce. The z9.6 model shows a nearly-symmetrical explosion, while the 9.0- M_\odot progenitor develops asymmetric explosions. The 10.0- M_\odot progenitor does not explode with the Baseline setup.

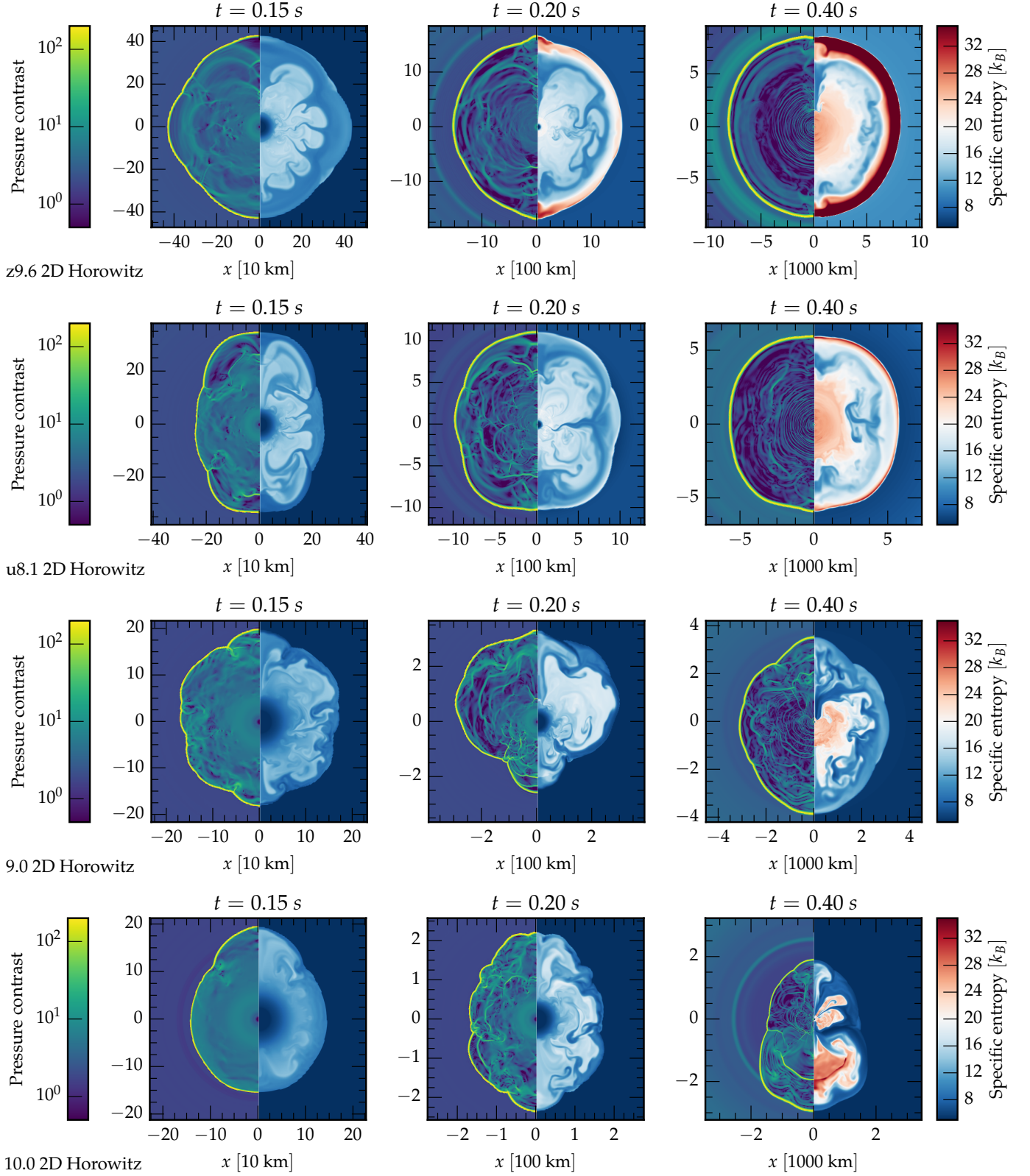


Figure 6. Entropy per baryon in k_B and pressure contrast ($r \cdot |\nabla p|/p$) profiles for the z9.6, u8.1, 9.0, and 10.0 progenitors evolved with many-body corrections at three representative times. This figure should be contrasted with Fig. 5. Compared to the Baseline setup, the inclusion of many-body corrections results in larger entropies and more violent convective overturn at early times. This is sufficient to turn the 10.0- M_\odot from a dud to a successful explosion. Many-body corrections also yield more symmetric explosions on for the 9.0- M_\odot progenitor.

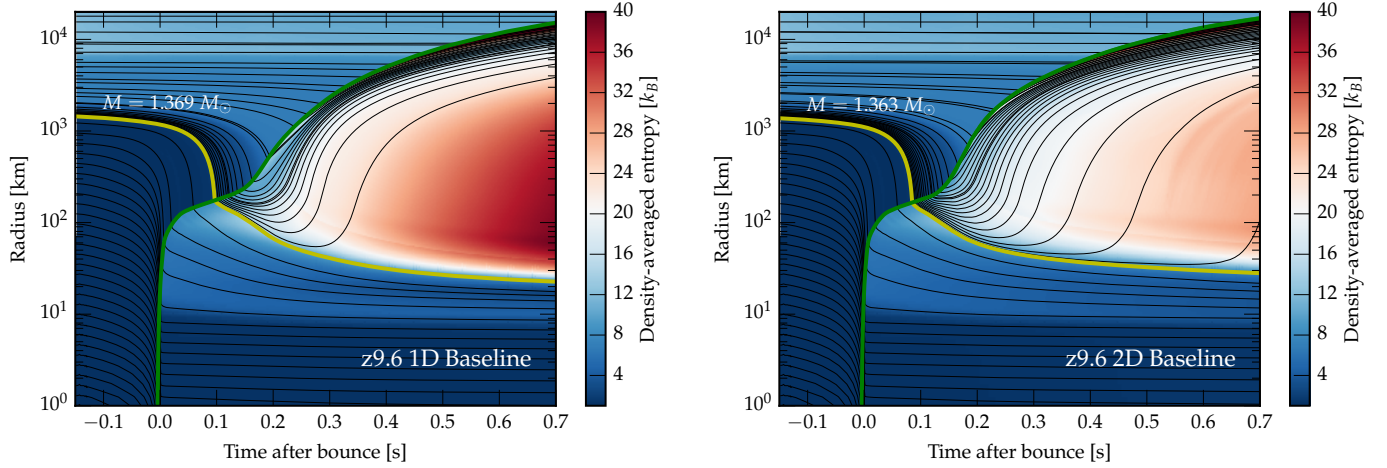


Figure 7. Evolution of the z9.6 progenitor in 1D (*left panel*) and 2D (*right panel*) with the Baseline setup. The green line denotes the average shock radius. The black lines are curves of constant enclosed baryonic mass (Lagrangian fluid elements in 1D). The yellow thick line denotes the final PNS mass cut. The curves are smoothed using a running average with a 5-ms window. The background color is the density-averaged entropy per baryon in k_B . 1D explosions generically result in the creation of low-density, high-entropy bubbles, which are smeared out by convection in 2D.

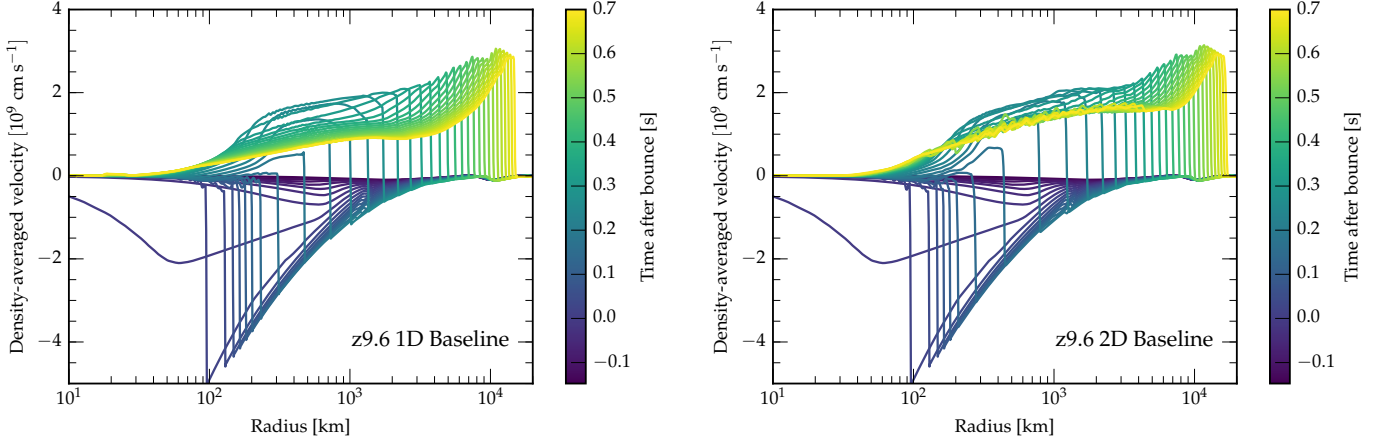


Figure 8. Density-averaged radial velocity in units of 10^9 cm s^{-1} for the z9.6 progenitor evolved with Baseline physics in 1D (*left panel*) and 2D (*right panel*). Multi-dimensional explosions result in larger velocities (and kinetic energies) behind the shock.

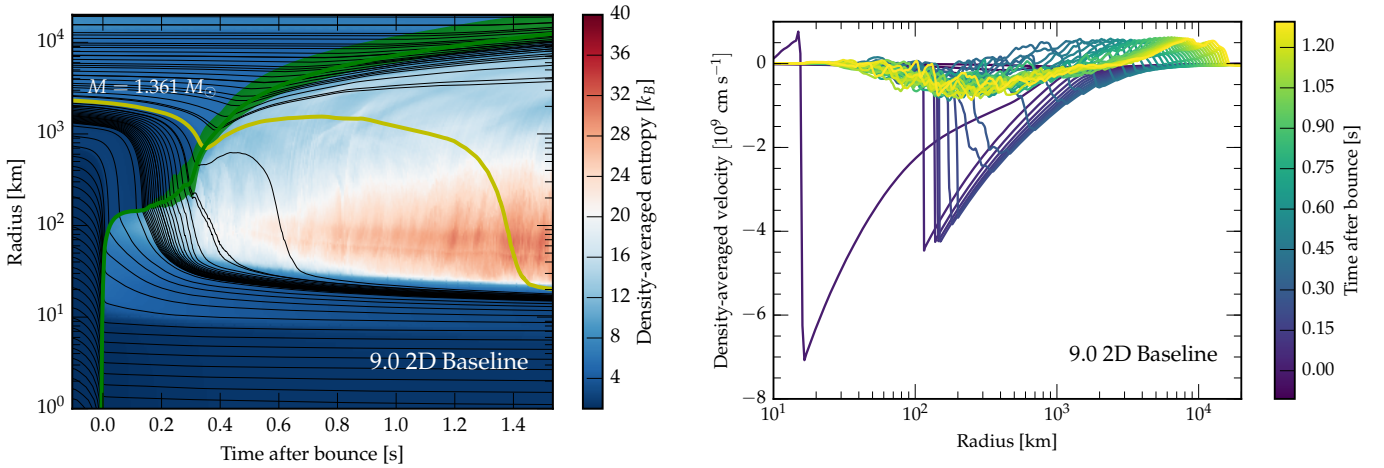


Figure 9. Evolution summary (*left panel*) and density-averaged velocity (*right panel*) for the 9.0 progenitor with the Baseline setup in 2D. The green shaded region in the left panel denotes the minimum and maximum shock radius. Curves in the left panel are smoothed using a running average with a 5-ms window. This model shows a marginal and asymmetric explosion. The velocities are positive behind the shock, signaling an overall expanding flow, but the expansion rate is much smaller than for the z9.6 progenitors (Fig. 8). Even when the shock reaches the outer boundary of our grid at 20,000 km, the velocity is still negative in regions behind the shock as a consequence of the partial fallback of the expanding plumes behind the shock.

Table 2
Summary of models and energy budget in the region $V : 100 \text{ km} \leq r \leq 20,000 \text{ km}$. Values are given at final simulation time.

Prog.	Setup	$10^3 \xi_{2.5}^a$	E_{bind}^b [10^{50} erg]	U^c [10^{50} erg]	E_0^d [10^{50} erg]	K_r^e [10^{50} erg]	K_θ^f [10^{50} erg]	E_g^g [10^{50} erg]	E_{tot}^h [10^{50} erg]	\dot{E}_{tot}^i [$10^{50} \text{ erg s}^{-1}$]	$E_{\nu_e}^j$ [10^{52} erg]	$E_{\nu_\mu}^k$ [10^{52} erg]	t_{end}^m [s]
n8.8	2D NoINS	-	0.000	0.307	-0.043	0.706	0.000	-0.079	0.892	0.103	1.465	3.202	0.701
n8.8	2D Baseline	-	0.000	0.426	-0.058	1.360	0.000	-0.112	1.616	0.190	1.526	2.740	0.663
n8.8	2D Horowitz	-	0.000	0.514	-0.084	2.260	0.000	-0.105	2.586	0.192	1.588	4.098	0.587
u8.1	2D NoINS	0.094	-0.018	0.415	-0.074	0.328	0.005	-0.140	0.535	0.068	1.967	1.504	1.121
u8.1	2D Baseline	0.094	-0.018	0.637	-0.093	0.762	0.002	-0.130	1.179	0.026	2.002	1.587	0.989
u8.1	2D Horowitz	0.094	-0.018	0.843	-0.120	1.184	0.002	-0.137	1.772	0.039	2.017	1.570	0.823
z9.6	2D NoINS	0.076	-0.008	0.326	-0.056	0.532	0.000	-0.085	0.717	0.024	1.714	1.253	0.839
z9.6	2D Baseline	0.076	-0.008	0.479	-0.070	1.065	0.001	-0.117	1.358	0.110	1.738	1.322	0.739
z9.6	2D Horowitz	0.076	-0.008	0.623	-0.099	1.765	0.001	-0.116	2.174	0.143	1.784	1.344	0.639
9.0	2D NoINS	0.038	-0.021	0.190	-0.047	0.080	0.009	-0.188	0.044	0.042	2.733	2.258	2.099
9.0	2D Baseline	0.038	-0.021	0.490	-0.112	0.181	0.011	-0.237	0.333	0.039	2.436	2.023	1.413
9.0	2D Perturb	0.038	-0.021	0.550	-0.110	0.259	0.014	-0.242	0.472	-0.039	2.502	2.096	1.560
9.0	2D Horowitz	0.038	-0.021	1.069	-0.185	0.868	0.008	-0.266	1.493	-0.013	2.302	1.856	1.097
10.0	2D NoINS	0.216	-0.095	0.309	-0.169	0.206	0.000	-0.637	-0.290	0.434	3.063	2.506	1.254
10.0	2D Baseline	0.216	-0.095	0.296	-0.163	0.198	0.000	-0.609	-0.277	0.347	3.270	2.759	1.298
10.0	2D Perturb	0.216	-0.095	0.938	-0.206	0.427	0.054	-0.803	0.409	0.573	3.309	2.836	1.615
10.0	2D Horowitz	0.216	-0.095	1.783	-0.389	0.670	0.057	-1.025	1.095	0.093	2.956	2.455	1.266
11.0	2D NoINS	7.669	-0.170	1.029	-0.552	0.432	0.000	-1.888	-0.979	0.891	3.187	2.612	1.262
11.0	2D Baseline	7.669	-0.170	2.269	-0.588	1.099	0.133	-1.916	0.997	0.381	3.344	2.857	1.569
11.0	2D Perturb	7.669	-0.170	2.128	-0.471	0.573	0.087	-1.987	0.331	1.946	3.642	3.165	2.046
11.0	2D Horowitz	7.669	-0.170	3.374	-0.848	1.090	0.149	-2.502	1.263	0.582	2.984	2.468	1.248

^aCompactness parameter.

^bEnvelope binding energy: $\int_{r>20,000 \text{ km}} \rho [\epsilon - GM/r] dV$.

^cTotal internal energy, excluding rest mass: $\int_V \rho \epsilon dV$.

^dTotal internal binding energy: $\int_V \rho \epsilon_0 dV$.

^eRadial kinetic energy: $0.5 \int_V \rho v_r^2 dV$.

^fNon-radial kinetic energy: $0.5 \int_V \rho (v_\theta)^2 dV$.

^gGravitational energy: $-\int_V \rho GM/r dV$.

^hTotal energy in the region of interest: internal, binding, kinetic, and gravitational.

ⁱRate of change of the net energy at the end of the simulation estimated from the last 10 ms of data.

^jTotal energy radiated in ν_e .

^kTotal energy radiated in ν_μ .

^lTotal energy radiated in " $\nu_{\mu'}$ ".

^mFinal simulation time (in seconds after bounce). This is the post-bounce time the maximum shock radius exceeds 19000 km for successful explosions.

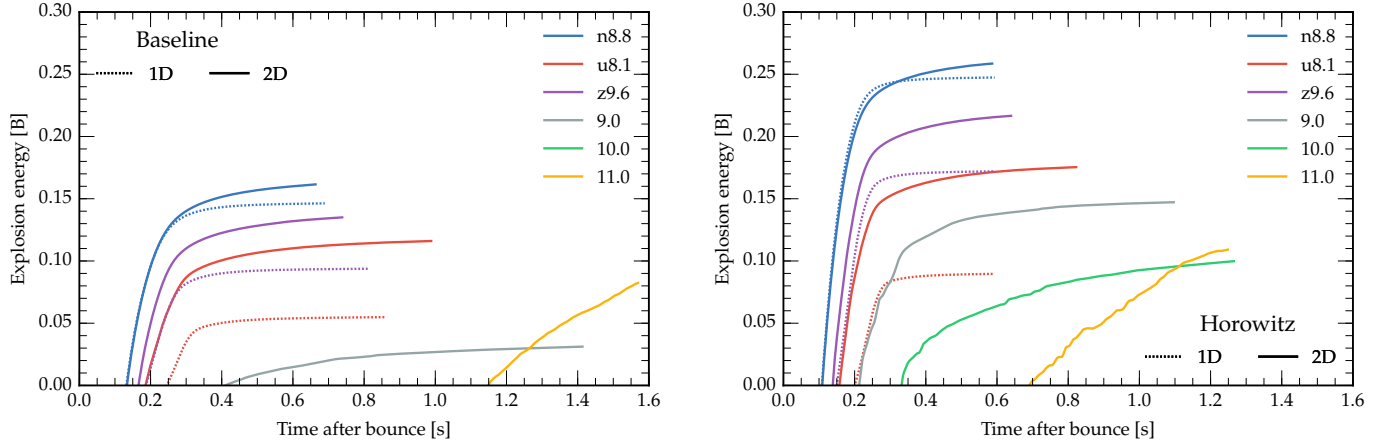


Figure 10. Net explosion energy in Bethe (10^{51} erg) in the region $100 \text{ km} \leq r \leq 20,000 \text{ km}$ for all progenitors in 1D and 2D with our Baseline setup (*left panel*) and with the inclusion of many-body corrections (*right panel*). We account for the binding energy of the envelope in the estimate of the explosion energy. The curves are smoothed using a running average with a 5-ms window. Many-body corrections have a tens of percent level impact on the explosion energies of all models, including those that explode in 1D.

the integrated total energy E_{tot} as the sum of internal U , kinetic K , and gravitational binding energy E_g . We also subtract the absolute value of the nuclear binding energy of the material E_0 . We remark that E_{tot} includes the contribution of the binding energy of the material exterior to the shock but interior to the computational domain. The values of these quantities at the end of our simulations are given in Table 2. There, we also quote the total energy liberated in each species of neutrinos E_{ν_e} , $E_{\bar{\nu}_e}$, and E_{ν_μ} . The final explosion energy can be estimated by subtracting in absolute value the binding energy of the material exterior to 20,000 km from E_{tot} . This is shown, as a function of time, in Fig. 10.

We find explosion energies ranging from a few percent of a Bethe ($1 \text{ B} \equiv 10^{51} \text{ erg}$), like the 9.0-Baseline 2D run, to values in excess of 1/4 of a Bethe, for the n8.8-Horowitz 2D simulation. The explosion energies we estimate for the z9.6 progenitor with either the Baseline or the Horowitz setup are, respectively, a factor of ~ 2 and ~ 3 larger than those reported by Melson et al. (2015b) for their 2D PROMETHEUS-VERTEX simulation. The estimated explosion energies for the z9.6, n8.8 and u8.1 in 2D are also similarly larger than those of COCONUT-VERTEX in 2D, as quoted in (Wanajo et al. 2017). Explosion energies for the progenitors from Sukhbold et al. (2016) have not been reported before, so no comparison is possible.

The difference is particularly striking when comparing the z9.6 progenitor in 1D. Melson et al. (2015b) find a shock runaway at much later times than in our 1D calculations and significantly smaller energies. Our 1D models are ~ 5 (~ 10) times more energetic with the Baseline (Horowitz) setup. These are significant differences considering that we are comparing 1D models for which the flow is necessarily laminar and the outcome should be deterministic. However, in Burrows et al. (2016), we showed that small differences in the input physics can result in diverging outcomes close to the threshold for explosion. This might explain the dissimilar outcomes. Indeed, even in 1D, the explosion energy of the z9.6 varies by a factor of ~ 5 between the Horowitz setup and the NoINS setup. This is still not enough to reconcile our results with those of Melson et al. (2015b), but should be

a motivation for a renewed effort in cross-code comparison and verification, in the spirit of Liebendoerfer et al. (2005).

The role of perturbations on the explosion energy (Table 2) is not completely clear. In the case of the 9.0- and 10.0- M_\odot progenitors, the inclusion of perturbations is beneficial. The first explodes with slightly larger explosion energy than without perturbations. The second goes from a failed to a successful, albeit underenergetic, explosion. Somewhat surprisingly, the 11.0-Perturb model has a smaller explosion energy than the unperturbed model. The reason is that 11.0-Perturb explosion entrains more bound mass and the ejecta lose more energy, while doing work on the infalling envelope of the star. On the other hand, note that the explosion energies for our the 11.0- M_\odot progenitor are still growing significantly at the end of our simulations, so the difference between the 11.0-Perturb and 11.0-Baseline models might be only transitory.

It is important to keep in mind that the explosion energies we quote are not final. Indeed the estimated explosion energy is still growing significantly at the time we stop the calculation for many of our simulations. This is not surprising in the light of the results of Müller (2015), who studied the development of an explosion in the 11.2- M_\odot progenitor from Woosley et al. (2002) and found the explosion energy to saturate only after several seconds. The inclusion of many-body corrections to the neutrino-nucleon scattering cross-sections in the Horowitz setup seems to be ameliorating this, since it results in earlier explosions that saturate more rapidly. On the other hand, the explosion energies saturate very rapidly for the n8.8, u8.1, and z9.6 progenitors and appear to have converged within the simulation time.

Another caveat is that our estimate of the explosion energy is more conservative than the commonly used “diagnostic explosion energy.” The former is computed as E_{tot} , but only integrated over unbound and/or radially expanding fluid elements (e.g., Buras et al. 2006; Müller et al. 2012b) and does not include the overburden of the material exterior to the shock. A comparison between the two is given in Fig. 11. There, we compute the diagnostic energy as the integral of the total energy density

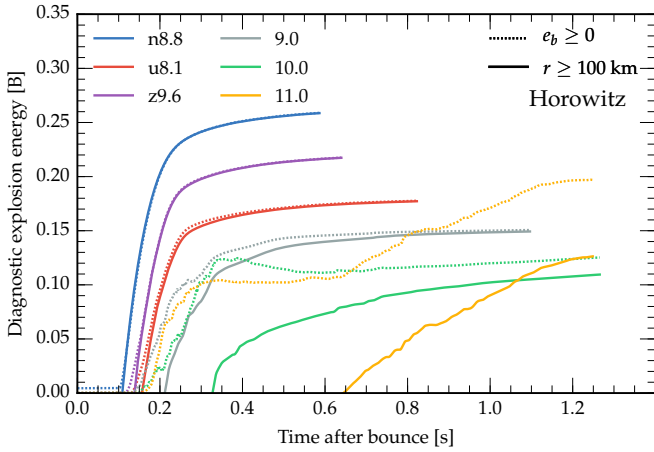


Figure 11. Diagnostic explosion energy E_{tot} in Bethe ($\equiv 10^{51}$ erg) integrated over the entire region where $r \geq 100$ km, or only for elements with positive net energy $e_b \geq 0$. The binding energy of the envelope has not been included in either calculation.

e minus the nuclear binding energy e_0 , $e_{\text{tot}} = e - e_0$, either integrated over regions where $e_{\text{tot}} \geq 0$, or over $100 \text{ km} \leq r \leq 20,000 \text{ km}$. For clarity, we did not include the binding energy of the envelope when computing E_{tot} in this plot, since its inclusion in the diagnostic explosion energy would be inconsistent. Beside this difference, the $r \geq 100 \text{ km}$ diagnostic energies in Fig. 11 are identical to the estimated explosion energies in Fig. 10. Obviously, the diagnostic energy integrated only over $e_{\text{tot}} \geq 0$ or $r \geq 100 \text{ km}$ should converge to the same value after a sufficiently long time. This is indeed the case for most of our models, especially the ECSN/ECSN-like explosions, where the explosion is close to spherically symmetric. However, significant differences persist until the end of our simulations for some progenitors. For example, in the 11.0-Horowitz run, the shock starts expanding ~ 200 ms after bounce, and some material becomes unbound. However, the energy behind the shock becomes sufficient to overcome the overburden only at later times, when the PNS wind becomes violent enough to create high-entropy bubbles behind the shock and the initially bound material at $r \geq 100 \text{ km}$ has accreted or has become unbound.

5. NEUTRINO RADIATION

We extract the properties of the neutrino radiation on a sphere placed at $10,000 \text{ km}$ from the center. Angle-averaged neutrino luminosities and rms neutrino energies are shown in Fig. 12. These are shown in the lab frame at infinity. Note that FORNAX evolves the fluid-frame neutrino-radiation moments and does not output the Eddington factor used for the evolution. For this reason, we convert the code output to the lab frame under the simplifying assumption of a forward-peaked radial neutrino distribution function. This assumption is not valid at the time the shock crosses $10,000 \text{ km}$ and results in small jumps which are particularly evident in the rms energies of heavy-lepton neutrinos for the n8.8 progenitor.

As in Burrows et al. (2016), we find the Horowitz setup to result in significantly higher neutrino luminosities and average energies in the first $\sim 0.2 \text{ s}$ after bounce. The heavy-lepton neutrinos are the most clearly affected,

since their opacity is dominated by neutrino-nucleon scattering. However, significant differences are seen for all species at early times. Over longer timescales, the differences are less pronounced. For the 10.0- and 11.0- M_{\odot} progenitors, which explode weakly (11.0- M_{\odot}) or not at all (10.0- M_{\odot}) with the Baseline setup, the neutrino luminosities after $\sim 0.3 \text{ s}$ are actually lower with the Horowitz setup in consequence of the smaller accretion luminosity. The averages energies for the 10.0- M_{\odot} model are also lower with the Horowitz setup than with the Baseline setup after $\sim 0.4 \text{ s}$. This is due to the successful explosion in the former simulation.

We quantify the degree of coupling between the neutrino radiation and the accretion flow in terms of the heating efficiency parameter η , defined as the ratio between the heating rate by neutrinos in the gain region, *i.e.*, the region bounded by the PNS and the shock with positive net neutrino heating, and the sum of the ν_e and $\bar{\nu}_e$ luminosities at infinity (*e.g.*, Marek & Janka 2009; Müller et al. 2012b,a). We show this quantity as a function of time in Fig. 13 for 1D and 2D models with the Baseline and the Horowitz setups. We recall that the same analysis was performed by Müller et al. (2012a) for the u8.1 progenitor. Their efficiency is slightly higher ($\sim 40\%$), but overall compatible with that in our Baseline 2D simulation. This suggests that the lower explosion energy they report is at least in part due to lower neutrino luminosities. Indeed, their ν_e and $\bar{\nu}_e$ luminosities are $\sim 20\%$ smaller than ours.

For the other models, the overall trend in η is that progenitors with larger accretion rates also show larger heating efficiencies. Spherically-symmetric (1D) simulations have significantly smaller heating efficiencies, with the possible exception of the n8.8 model, which is very close to being spherically symmetric even in 2D. The increased heating efficiency in 2D is in part due to the longer dwelling time of material in the gain region (Burrows et al. 1995; Murphy & Burrows 2008; Dolence et al. 2013) or, equivalently (Müller et al. 2012b) to the growth of the mass in the gain region.

The many-body corrections implemented in the Horowitz setup result in a significant increase of the heating efficiency, in some cases up to $\sim 50\%$. At least at early times, before the evolutionary paths of the Baseline and Horowitz simulations start to diverge, this improvement can be attributed to the hardening of the neutrino spectra with the Horowitz setup, which results in a more tight coupling with the material. After the explosions set in, the differences between the Baseline and Horowitz efficiencies are in good part due to the fact that the Horowitz explosions are more spherical and entrain more mass.

6. PROTONEUTRON STARS

As commonly done in the CCSN-mechanism literature, we define as PNS radius the radius at which the angle-averaged density is $10^{11} \text{ g cm}^{-3}$. We monitor PNS radii and the baryonic mass they enclose to estimate the final remnant radius. These quantities are shown in Figs. 14 and 15. PNS masses and accretion rates at the end of our simulations are also given in Table 3. There we also quote the corresponding gravitational mass for a cold, deleptonized NS estimated using the approximate fit of Timmes et al. (1996).

The PNS masses, as was the case for the explosion

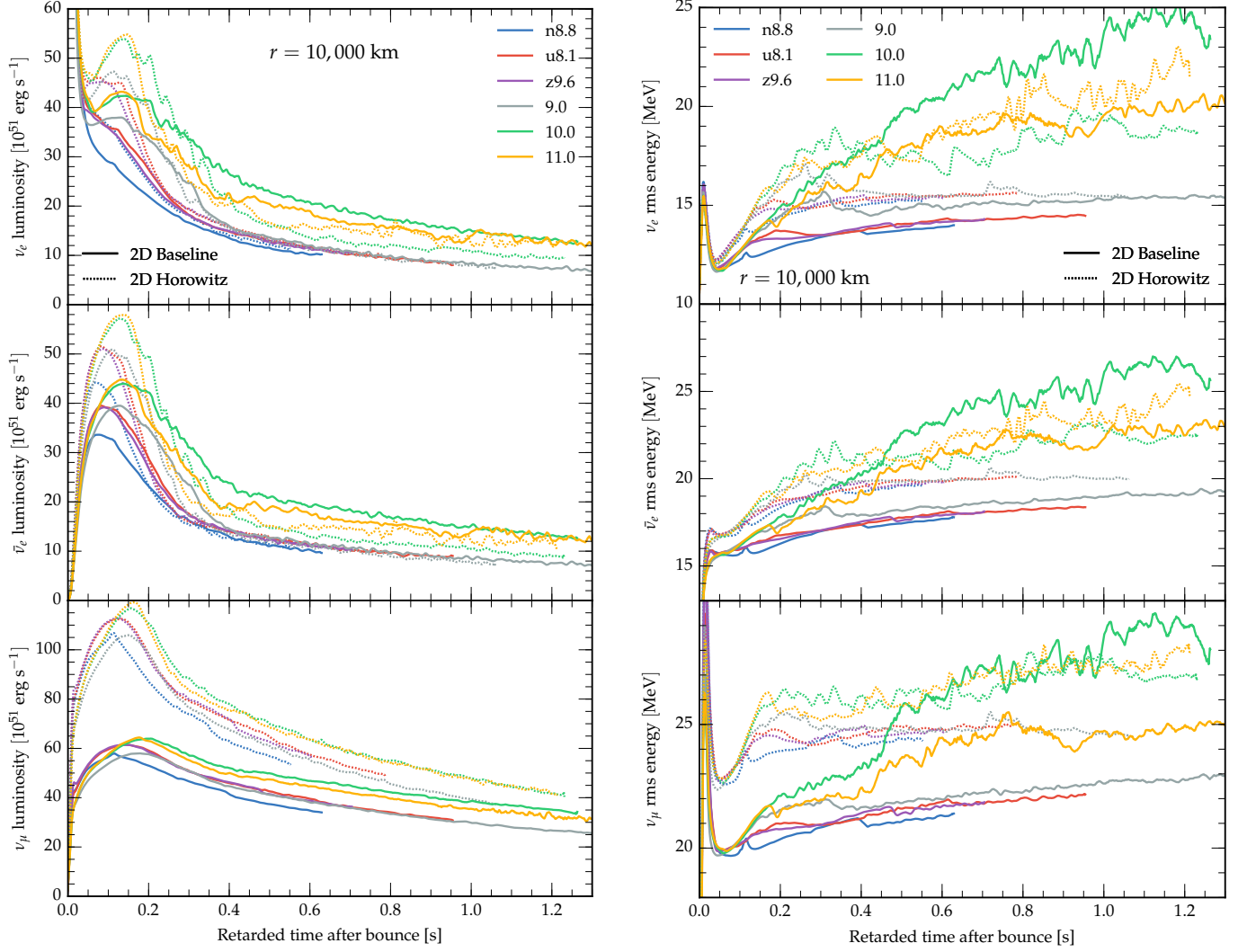


Figure 12. Neutrino luminosity (*left panel*) and rms energies (*right panel*) at 10,000 km as a function of the retarded time. Here, ν_μ denotes the sum of all heavy-lepton neutrino species. The curves are smoothed using a running average with a 5-ms window. As was the case with more massive progenitors (Burrows et al. 2016), many-body corrections increase neutrino luminosities and average energies.

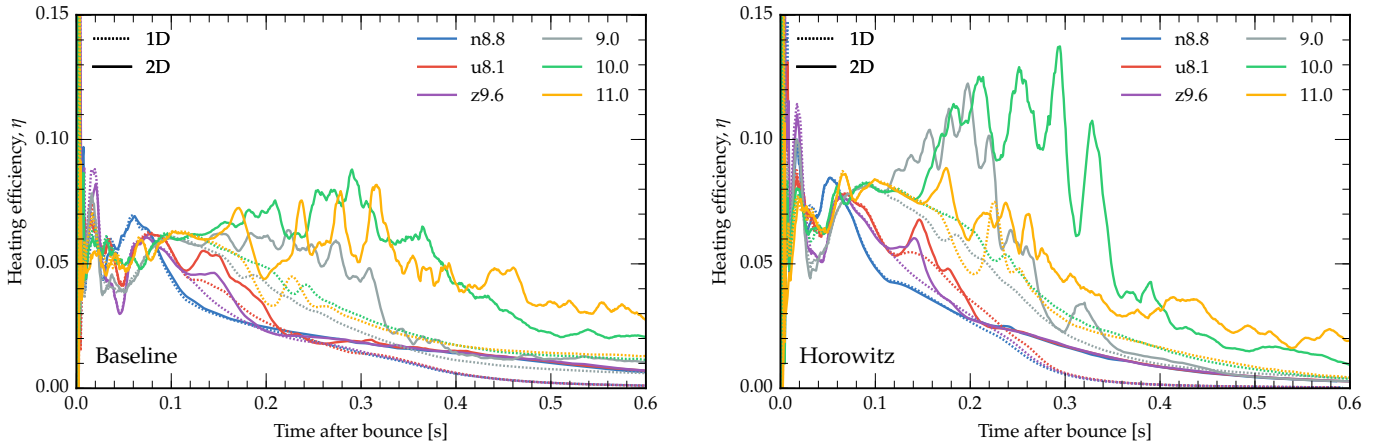


Figure 13. Heating efficiency, η , for 1D and 2D models with the Baseline setup (*left panel*) and with many-body corrections included in the neutrino opacities (*right panel*). Many-body effects result in a hardening of the neutrino radiation which, in turn, leads to a better coupling of ν_e and $\bar{\nu}_e$ with the material in the gain region.

Table 3
PNS star masses at the final simulation time.

Prog.	Setup	$M_{\text{Baryon}}^{\text{a}}$ [M_{\odot}]	$M_{\text{Grav}}^{\text{b}}$ [M_{\odot}]	$\dot{M}_{\text{Baryon}}^{\text{c}}$ [$M_{\odot} \text{ s}^{-1}$]
n8.8	1D NoINS	1.304	1.197	0.002
n8.8	1D Baseline	1.297	1.191	0.003
n8.8	1D Horowitz	1.294	1.188	0.000
n8.8	2D NoINS	1.301	1.194	0.005
n8.8	2D Baseline	1.294	1.188	0.007
n8.8	2D Horowitz	1.291	1.186	0.005
u8.1	1D NoINS	1.392	1.271	0.009
u8.1	1D Baseline	1.378	1.259	0.002
u8.1	1D Horowitz	1.376	1.258	0.002
u8.1	2D NoINS	1.377	1.258	0.001
u8.1	2D Baseline	1.369	1.252	0.002
u8.1	2D Horowitz	1.366	1.249	0.001
z9.6	1D NoINS	1.376	1.257	0.001
z9.6	1D Baseline	1.369	1.252	0.001
z9.6	1D Horowitz	1.365	1.248	0.002
z9.6	2D NoINS	1.370	1.252	0.003
z9.6	2D Baseline	1.364	1.247	0.005
z9.6	2D Horowitz	1.360	1.244	0.003
9.0	1D NoINS	1.373	1.255	0.017
9.0	1D Baseline	1.373	1.255	0.016
9.0	1D Horowitz	1.369	1.251	0.024
9.0	2D NoINS	1.378	1.259	0.007
9.0	2D Baseline	1.361	1.244	0.002
9.0	2D Perturb	1.357	1.242	0.002
9.0	2D Horowitz	1.336	1.224	0.000
10.0	1D NoINS	1.518	1.376	0.074
10.0	1D Baseline	1.516	1.374	0.077
10.0	1D Horowitz	1.498	1.359	0.097
10.0	2D NoINS	1.538	1.393	0.045
10.0	2D Baseline	1.540	1.394	0.040
10.0	2D Perturb	1.516	1.374	0.023
10.0	2D Horowitz	1.455	1.324	0.011
11.0	1D NoINS	1.526	1.383	0.095
11.0	1D Baseline	1.511	1.370	0.105
11.0	1D Horowitz	1.508	1.368	0.111
11.0	2D NoINS	1.550	1.402	0.088
11.0	2D Baseline	1.517	1.375	0.046
11.0	2D Perturb	1.509	1.369	0.027
11.0	2D Horowitz	1.455	1.324	0.039

^aPNS baryonic mass.

^bPNS gravitational mass.

^cPNS accretion rate.

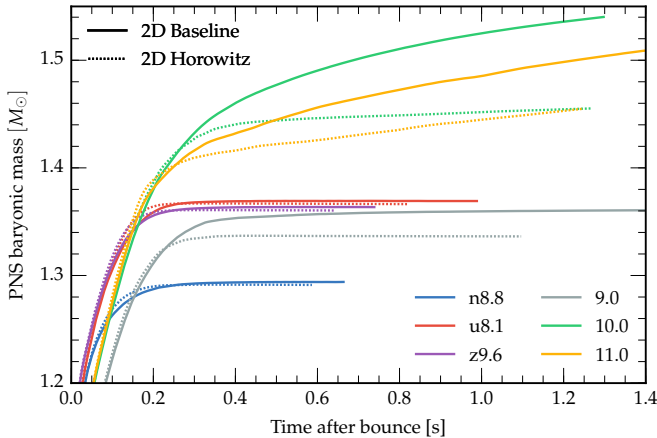


Figure 14. PNS baryonic masses for the 2D simulations with the Baseline and Horowitz setups. The curves are smoothed using a running average with a 5-ms window. Many-body corrections to the neutrino-nucleon scattering opacities result in earlier, more vigorous explosions and, consequently, smaller PNS masses.

energy, are still not converged for the 11.0- M_{\odot} progenitor at the end of our simulations, although the Horowitz run appears to be close to convergence. The PNS mass for the 10.0- M_{\odot} model is obviously only converged with the Horowitz and Perturb setup, which explode, while black-hole formation appears inevitable for the other setups.

Notwithstanding these caveats, we find that all our progenitors produce PNSs with gravitational masses below 1.4 M_{\odot} . The n8.8 progenitor produces PNSs with gravitational masses as low as 1.186 M_{\odot} , suggesting that an ECSN/ECSN-like explosion might be able to explain the origin of the low-mass companion in the double NS system J0453+1559. This has been recently measured, using the advance of periastron and the Shapiro delay, to have a mass of $1.174 \pm 0.004 M_{\odot}$ (Martinez et al. 2015). This scenario is also plausible in light of the study by Tauris et al. (2015), who showed that an ECSN is a possible outcome of the evolution of ultra-stripped metal cores in tight binaries, such as those producing relativistic double-NS systems like J0453+1559.

We find the PNS radii (Fig. 15) to follow tracks that are largely independent of the progenitor or the PNS mass, as do Bruenn et al. (2016) and Summa et al. (2016). The reason is that the density drops sharply at the surface of the PNS so that the ambient pressure has a negligible influence on the structure of the central object. This is, rather, determined by the competition between its internal pressure and gravity. Instead, the radii are sensitive to changes in the microphysical treatment, which determines the rate at which the PNS deleptonizes and loses thermal support and to the dimensionality (1D vs. 2D). The impact of the microphysics is easily understood from the fact that the contraction of the PNS is mostly set by the rate of deleptonization and core cooling. These in turn depend in the first second after bounce on the neutrino opacity of matter at densities between 10^{11} and $10^{13} \text{ g cm}^{-3}$. For instance, the many-body corrections included in the Horowitz setup result in a faster deleptonization and contraction of the PNS.

The reason for the faster PNS contraction in 1D is more easily understood considering the z9.6 progenitor with the Baseline setup, which explodes both in 1D and in 2D. Its PNS evolution is shown in Fig. 16. In the first few hundreds of milliseconds after bounce, the convection inside the PNS is buried deep below the surface and its impact on supernova evolution is limited, as has been documented in detail by Buras et al. (2006) and Dessart et al. (2006). However, over timescales longer than those considered in either of those works, starting from ~ 0.25 s after bounce, the surface of the PNS shrinks to the point of entering in contact with the inner PNS convection, which then becomes dynamically important. This can be seen in Fig. 16, where the region affected by the inner convection is identifiable by its small radial entropy gradient with entropy per baryon evolving from $\sim 4.5 \text{ k_B}$ to $\sim 3 \text{ k_B}$ as the PNS cools down.

Starting from this moment, the 1D and 2D evolutions begin to diverge. In 1D, the neutrino cooling of the surface is not compensated by convection and leads to an increasingly steep entropy inversion. The pressure support in the exterior layers of the PNS drops rapidly and leads to an increased compactness, with respect to the 2D evolution, of the regions with densities between 10^{11} and $10^{13} \text{ g cm}^{-3}$. These regions are instead inflated in 2D

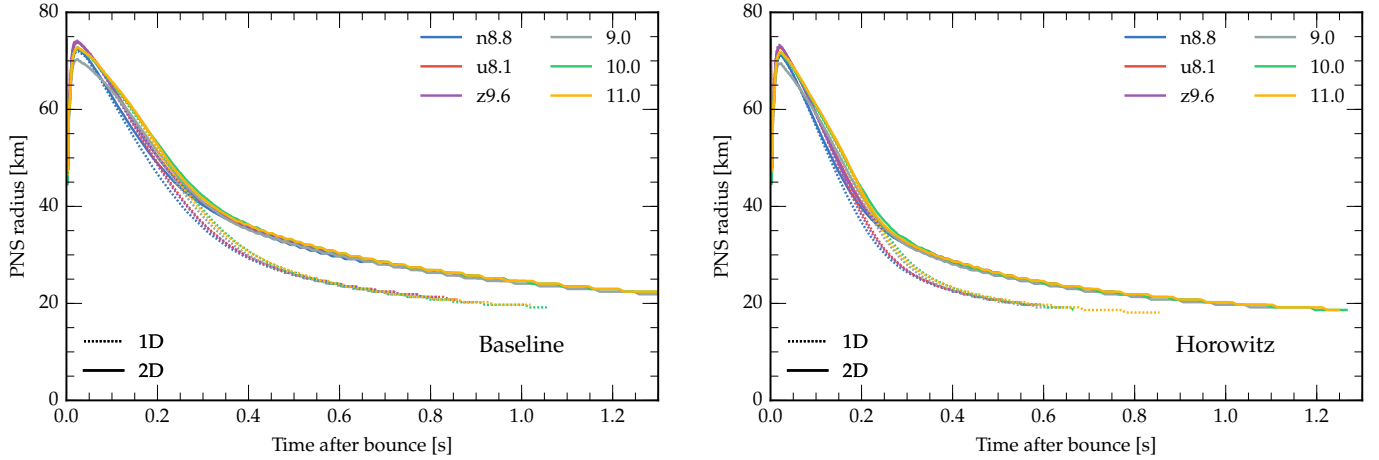


Figure 15. PNS radii in 1D and 2D with Baseline physics (*left panel*) and with many body-corrections (*right panel*). Curves are smoothed using a running average with a 5-ms window. Many-body corrections result in faster PNS contraction rates, but the largest differences are between 1D and 2D simulations. One-dimensional models predict faster contraction rates of the PNS starting from ~ 0.25 s after bounce.

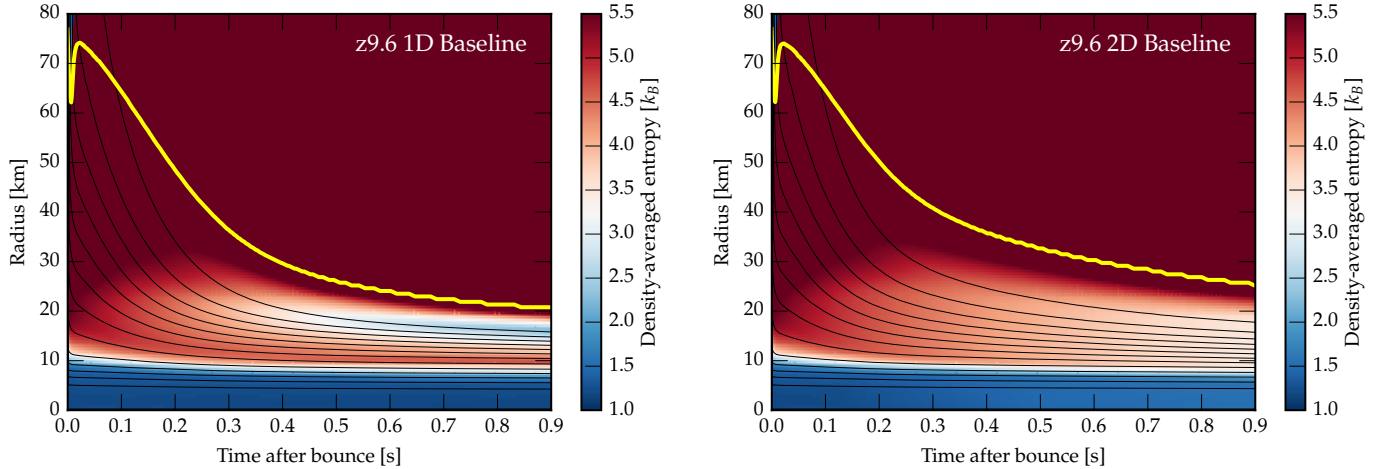


Figure 16. Evolution of the PNS for the z9p6 progenitor in 1D (*left panel*) and 2D (*right panel*) with the Baseline setup. The black lines are curves of constant enclosed baryonic mass. The yellow thick line denotes the PNS radius. The curves are smoothed using a running average with a 5-ms window. The background color is the density-averaged entropy per baryon in k_B . The PNS radius contracts to the point of touching the inner convection region, visible as the almost constant averaged entropy region exterior to ~ 10 km in 2D, at ~ 0.4 s after bounce. As a consequence, the subsequent evolution of the PNS is drastically different in 1D and 2D.

by the deposition of entropy and lepton number due to convective transport. The structure of the layers below the inner convection region is also affected, with the core of the PNS reaching higher densities and compactness in 1D.

PNS convection also leaves a strong imprint on the neutrino luminosities, which is boosted by up to a factor ~ 2 at late times, as can be seen from Fig. 17. This seems to be the main reason for the enhanced growth of the explosion energy for the n8.8, u8.1, and z9.6 progenitors in 2D at late times. While this accounts only for a relatively small fraction of the explosion energy (see Fig. 10), the role of PNS convection might be much more important for more massive progenitors that explode later in multi-dimensional simulations.

We remark that O’Connor & Couch (2015) also reported modest increases in the heavy-lepton neutrino luminosities due to PNS convection. However, since they considered models that did not explode in 1D, they might have underestimated the effect of convection, since the $\bar{\nu}_e$ and, in particular, the ν_e luminosities are significantly

affected by accretion. Indeed, while the n8.8 and u8.1 show very similar differences between 1D and 2D as does the z9.6, this is not the case for the 9.0, 10.0, and 11.0, which show more similar luminosities in 1D and 2D.

Müller & Janka (2014) also considered the z9.6 progenitors over a long timescale and found luminosities very close to ours with the Baseline setup. Their luminosity was $L_{\nu_e} \sim 10^{52} \text{ erg s}^{-1}$ at 0.6 s after bounce, the same value we also find (Fig. 17), but they did not present a comparison with the corresponding 1D evolution. The n8.8 was considered in 1D by Fischer et al. (2010), who found essentially the same luminosity as in our 1D n8.8-Baseline model ($L_{\nu_e} \simeq 6 \cdot 10^{52} \text{ erg s}^{-1}$ at 0.6 s after bounce). This is, however, a factor ~ 2 smaller than in our 2D calculations for the n8.8. These considerations are all additional indirect confirmation that the impact of PNS convection has been underestimated.

7. CONCLUSIONS

We have revisited the explosion of low-mass iron-core SNe and O-Ne-Mg gravitational-collapse SNe with a new set of neutrino-radiation hydrodynamics simulations in

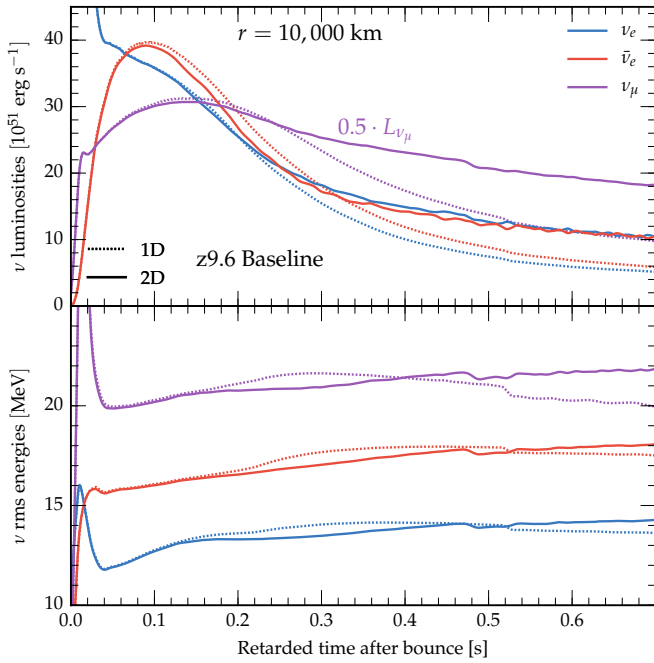


Figure 17. Neutrino luminosity (*top panel*) and rms energies (*bottom panel*) at 10,000 km as a function of the retarded time for the z9.6 progenitor evolved with the Baseline setup. Here, ν_μ denotes the sum of all heavy-lepton neutrino species and their associated luminosity. It has been rescaled by a factor 0.5 to improve the readability of the plot. The curves are smoothed using a running average with a 5-ms window. Neutrino-driven convection below the neutrinospheres result in a boost of the luminosity compared to 1D models.

1D (spherical symmetry) and in 2D (axial symmetry). Our simulations included the effects of general relativity in an approximate way and state-of-the-art neutrino transport and weak reactions. Of the six progenitors we have considered, one, the n8.8 from Nomoto (1984, 1987), is the prototype of an ECSN. Two, the $10^{-4} Z_\odot$ and zero metallicity u8.1 and z9.6 from Müller et al. (2012a) and Müller et al. (2013), have iron cores, but a structure similar to that of the n8.8. The 9.0-, 10.0-, 11.0- M_\odot solar-metallicity progenitors from Sukhbold et al. (2016) share some similarities with the n8.8, but are overall closer to the “canonical” CCSN progenitors considered in the past in the CCSN-mechanism literature.

As in previous studies (Kitaura et al. 2006; Janka et al. 2008; Burrows et al. 2007; Fischer et al. 2010; Müller et al. 2012a; Janka et al. 2012; Müller et al. 2013; Müller & Janka 2014; Melson et al. 2015b; Wanajo et al. 2017), we find that the n8.8, u8.1, and z9.6 progenitors typically explode easily, even in 1D. On the other hand, the low-mass, but solar-metallicity, progenitors with iron cores from Sukhbold et al. (2016) do not explode in 1D and, in some cases, like the 10.0 with our Baseline setup, fail to explode also in 2D. Taken together with our previous findings (Burrows et al. 2016), our results show that solar-metallicity iron-core SNe do not explode in 1D and are not even necessarily easier to explode than higher-mass stars.

The failure to explode of the 10.0- M_\odot progenitor from Sukhbold et al. (2016) reported here, as well as the failed explosions of the 12 and 15- M_\odot progenitors from Woosley & Heger (2007) we reported in Burrows et al. (2016) with Baseline setup, is surprising in the light of

the successful explosion of other progenitors, like the 9.0- and 11.0- M_\odot progenitors from Sukhbold et al. (2016) and the 20- and 25- M_\odot progenitors from Woosley & Heger (2007), with the same setup. Our findings are in tension with explodability criteria related to the ZAMS mass (Heger et al. 2003) or to the progenitor compactness and related parameters (O’Connor & Ott 2011, 2013; Ugliano et al. 2016; Nakamura et al. 2015; Pejcha & Thompson 2015; Ertl et al. 2016). Other circumstantial evidence of limitations in the existing explodability criteria is the order in which explosions develop in Summa et al. (2016) and the results of O’Connor & Couch (2015). The latter also considered the 12-, 15-, 20-, and 25- M_\odot progenitors from Woosley & Heger (2007) and found explosions in approximate GR for all progenitors except the 12 M_\odot . Taken together, all of these results suggest that, while some properties of the explosions are correlated with the compactness of the progenitor (e.g., O’Connor & Ott 2013; Nakamura et al. 2015), the explodability is not. Whether an explosion is successful or not depends on a competition between accretion and neutrino heating (Burrows & Goshy 1993; Janka 2000; Suwa et al. 2016; Murphy & Dolence 2017; Gabay et al. 2015) which, in our opinion, has yet to be expressed in terms of the progenitor properties in a satisfactory way.

We have systematically studied the effect of perturbations and of changes in the treatment of neutrino-matter interactions, with emphasis on the impact of the many-body corrections to the neutrino-nucleon scattering cross section derived by Horowitz et al. (2016). As was the case for the higher ZAMS mass progenitors (Burrows et al. 2016), we have found that relatively small changes are amplified by the proximity to the threshold for explosion and can lead to qualitatively different outcomes. For instance, the 10.0- M_\odot model turns from a dud into an explosion with the inclusion of perturbations or with many-body corrections. This sensitivity to initial conditions and/or physics setup applies also in 1D to those progenitors that are sufficiently close to the threshold for explosion. For example, the u8.1 fails to explode in 1D without the inclusion of inelastic scattering on nucleons, but succeeds when those are included. The reason for the diverging outcomes with different microphysical suites is easily understood in terms of the neutrino-radiation intensity and hardness, which directly translate into the efficiency of the energy deposition by neutrinos.

We have estimated explosion energies by following the development of the explosions over long timescales and until the shock has reached 19,000 km in 2D simulations. While the explosion energy for the 11.0- M_\odot progenitor is still far from saturated, for the others we have found saturated explosion energies, computed accounting for the binding energy of the entire star, of a fraction of a Bethe. These values are in the expected range for low-mass progenitors (Utrobin & Chugai 2013; Spiro et al. 2014; Sukhbold et al. 2016). We remark that, while the ECSN/ECSN-like explosions are nearly-spherical and we do not expect their explosion energies to change significantly in 3D (Melson et al. 2015b), it is likely that the asymmetric explosions we observe for the 9.0-, 10.0-, and 11.0- M_\odot progenitors will be quantitatively different in 3D (Müller 2015).

For the progenitors that have been evolved in the past by other groups (the n8.8, u8.1, and z9.6) we obtain ex-

plosion energies in 2D that are typically a factor ~ 2 larger than those of the Garching group (Janka et al. 2008; Melson et al. 2015b; Wanajo et al. 2017). The differences are even larger for the z9.6 in 1D, for which we predict energies between ~ 2 and ~ 10 times larger than theirs, depending upon which microphysics suite we use. The diversity in the explosion energies is particularly troubling in 1D, where the absence of turbulence should make the evolution deterministic. Nevertheless, the fact that our estimates change significantly as a result of relatively small modification to the neutrino opacities suggests that our results and those of Melson et al. (2015b) could probably be reconciled by tweaking our microphysical treatment. We did not attempt it here, but the use of models that explode in self-consistent 1D simulations appears promising in the context of studies on the impact of the microphysics on the explosion mechanism, or within a renewed effort to cross-validate CCSN codes. It seems that, more than ten years after Lieben-derfer et al. (2005), there is still a good case for more 1D (and ostensibly 2D) comparisons. This is an avenue we will pursue in the near future.

For exploding models, we have found final PNS masses that have reached saturation within the simulation time, with the exception of the 11.0 progenitor. We find that ECSNe can explain the low-mass tail of the observed NS mass distribution. The n8.8 progenitor with the Horowitz setup leaves behind a NS with baryonic (gravitational) mass of $1.291 M_{\odot}$ ($1.186 M_{\odot}$), very close to the lowest accurately measured NS gravitational mass of $1.174 \pm 0.004 M_{\odot}$ (Martinez et al. 2015).

We studied the evolution of the PNS, focusing on those progenitors that explode both in 1D and 2D. These have nearly identical boundary conditions, allowing us to quantify the role of multi-dimensional effects on the long-term evolution of the PNS. We have found that the PNS contraction rate slows down significantly in 2D compared to 1D, starting from ~ 0.25 s after bounce. At this time, the PNS surface has contracted sufficiently to enter into contact with the inner PNS convection region. The transport of lepton number and thermal energy by the PNS convection then inflates the region with densities between 10^{11} and 10^{13} g cm $^{-3}$, causing the decrease of the contraction rate.

We have also found PNS convection to be responsible for a boost of the neutrino luminosities for all species at late times ($\gtrsim 0.5$ s after bounce) by a factor ~ 2 . This contributed only a small percentage of the explosion energy for the ECSN/ECSN-like progenitor for which this effect is easily quantifiable. However, PNS convection is likely to be more important for massive progenitors that explode late, when the PNS surface has already receded sufficiently close to the PNS convection region. Our results, together with pieces of evidence from Fischer et al. (2010) and Müller & Janka (2014), strongly suggest that the impact of PNS convection has been underestimated in the past (Buras et al. 2006; Dessart et al. 2006). Our findings provide an additional reason, beside the need to account for continued accretion at late times (Müller & Janka 2014), for the importance of multi-D simulations in the modeling of the early neutrino signal from cooling PNSs (e.g., Fischer et al. 2010; Hudepohl et al. 2010; Roberts 2012; Nakazato et al. 2013; Roberts & Reddy 2016).

The main limitation of this work is the assumption of axisymmetry. This has been necessary given the large computational cost of 3D simulations with state-of-the-art microphysics, which prevents a systematic study with multiple progenitors, as is the present one. This is attested by the scarcity of 3D simulations with full microphysics (Hanke et al. 2013; Tamborra et al. 2014; Melson et al. 2015b,a; Lentz et al. 2015). However, moving to 3D will ultimately be required and will be a goal of our future work.

ACKNOWLEDGMENTS

The authors acknowledge Chuck Horowitz, Evan O'Connor, and Todd Thompson for productive conversations concerning, insight into, and help with the microphysics, and Ernazar Abdikamalov, Sean M. Couch, Luke F. Roberts and Christian D. Ott for fruitful discussion on the nature of core-collapse supernovae. Support was provided by the Max-Planck/Princeton Center (MPPC) for Plasma Physics (NSF PHY-1144374) and the NSF PetaApps program, under award OCI-0905046 via a subaward no. 44592 from Louisiana State University to Princeton University. DR gratefully acknowledges support from the Schmidt Fellowship. JD acknowledges support from a Laboratory Directed Research and Development Early Career Research award at LANL. The authors employed computational resources provided by the TIGRESS high performance computer center at Princeton University, which is jointly supported by the Princeton Institute for Computational Science and Engineering (PICSciE) and the Princeton University Office of Information Technology and by the National Energy Research Scientific Computing Center (NERSC), which is supported by the Office of Science of the US Department of Energy (DOE) under contract DE-AC03-76SF00098. The authors express their gratitude to Ted Barnes of the DOE Office of Nuclear Physics for facilitating their use of NERSC. This paper has been assigned a LANL preprint # LA-UR-17-20973.

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